Charm and hidden charm scalar mesons in the nuclear medium

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Abstract

We study the renormalization of the properties of low lying charm and hidden charm scalar mesons in a nuclear medium, concretely of the $D_{s0}(2317)$ and the theoretical hidden charm state $X(3700)$. We find that for the $D_{s0}(2317)$, with negligible width at zero density, the width becomes about 100 MeV at normal nuclear matter density, while in the case of the $X(3700)$ the width becomes as large as 200 MeV. We discuss the origin of this new width and trace it to reactions occurring in the nucleus, while offering a guideline for future experiments testing these changes. We also show how those medium modifications will bring valuable information on the nature of the scalar resonances and the mechanisms of the interaction of D mesons with nucleons and nuclei.

1 Introduction

The modification of the properties of elementary particles in nuclei is a rich field that helps simultaneously to learn about excitation mechanisms in the nucleus as well as properties of the elementary particles [1]. Important features about the nature of certain particles can be better observed in nuclei. Take for instance the $\Lambda(1520)$. This resonance couples strongly to $\pi\Sigma(1385)$, but this is hardly visible in the decay of the $\Lambda(1520)$ particle since there is no phase space for it. When this resonance is placed inside a nucleus the pion can become a particle hole excitation (ph) and one gains 140 MeV of phase space for the decay, which results in a considerable increase in the width of the $\Lambda(1520)$ in the medium [2]. Take another example: the ω meson. This meson decays into 3π, which is supposed to go through $\rho\pi$, but there is no

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phase space for decay in this channel except for the ρ width. Once again the $\pi\pi$ decay channel of the ρ will be modified in the nucleus, as well as the individual π, which can become again a ph excitation producing a much larger phase space for the decay of the ω. It is then spectacular to find that the width of the ω in the medium at nuclear matter density becomes as large as $100 - 150$ MeV [3, 4]. This discussion is only qualitative and more subtle details must be considered as done in other works [5, 6, 7].

Among so many examples in the literature, the case of the renormalization of the properties of the scalar mesons has played a special role. One reason is that there is a long debate on the nature of the scalar mesons, with different assumptions about their nature as $q\bar{q}$ states, $K\bar{K}$ molecules, mixtures of $q\bar{q}$ with meson-meson components, or dynamically generated resonances from the interaction of coupled channels of two pseudoscalars [8]-[26]. We will study here the properties in the medium of two dynamically generated states, meaning states made out of two mesons. Pioneer work in the study of dynamically generated resonances was done in [8, 9], where starting from one seed of $q\bar{q}$, a large meson cloud was demanded to explain data of the low lying scalars. Further work in this direction followed using the unitary coupled channel chiral approach in [24, 25, 26]. In this sense, the predictions of the different models on the medium modification of these resonances are important in order to put stringent constraints on the different assumptions about the nature of the resonances.

Among the low lying scalars, the $\sigma(600)$ has been the most studied. Several theoretical approaches have predicted strong medium effects on the $\pi\pi$ interaction in the scalar isoscalar (σ) channel. In [27], the σ and π mesons at high density are studied within the Nambu-Jona-Lasinio model and the authors find that the m_σ drops subtantially with the density, whereas the m_π increase at higher densities. The π was previously studied at finite temperature and density within the same model in [28]. In Ref. [29], Hatsuda et al. studied the σ propagator in the linear σ model and found an enhanced and narrow spectral function near the 2π threshold caused by the partial restoration of the chiral symmetry, where m_σ would approach m_π. The same conclusions were reached using the nonlinear chiral Lagrangians in Ref. [30]. Similar results, with large enhancements in the $\pi\pi$ amplitude around the 2π threshold, have been found in quite different approaches by studying the $s-$wave, $I = 0$ $\pi\pi$ correlations in nuclear matter [31]. In these cases the modifications of the σ channel are induced by the strong $p-$wave coupling of the pions to the particle-hole (ph) and Δ-hole (Δh) nuclear excitations. A more recent and updated theoretical work on the evolution of the σ poles in the medium can be seen in [32].

On the experimental side, there are also several results showing strong medium effects in the σ channel at low invariant masses in the $A(\pi, 2\pi)$ [33] and $A(\gamma, 2\pi)$ [34] reactions, which have been the object of study in [35], [36] and [37].

The $f_0(980)$ and $a_0(980)$ scalar mesons have also been analyzed in the nuclear medium [38], but unlike the case of the σ, no experimental action has been taken in this case.

In the present paper we retake research along these lines and study the medium modification of the scalar mesons which are dynamically generated in the charm sector. Concretely, we shall study the medium modification of the $D_{s0}(2317)$, which is

obtained in the theoretical studies of [39]-[42]. Within the context of lattice gauge theories, in [43] the authors find hints that there is a DK bound state that can be identified with the $D_{s0}(2317)$. In addition, we shall also study the medium modifications of a hidden charm scalar meson predicted in [42], for which there are some indications that could have been observed in the experiment of the Belle collaboration [44] through the reanalysis done in [45]. In this experiment a broad bump is seen in the $D\bar{D}$ mass distribution around the $D\bar{D}$ threshold, which in [45] was shown to be better explained by the $X(3700)$ pole of the $D\bar{D}$ amplitude below threshold than by the Breit Wigner distribution proposed in the experimental paper. However, it should be noted that a better fit to the data with a Breit Wigner distribution than the one of [44] was obtained in [45], yet with a χ^2 value slightly worse than the one obtained with the pole below threshold.

2   Brief discussion on the dynamical generation of the $D_{s0}(2317)$ and $X(3700)$

We follow the details of [42], where a Lagrangian is taken for the interaction of two pseudoscalar mesons. The Lagrangian is an extrapolation to $SU(4)$ of the $SU(3)$ chiral Lagrangian used in [24, 26] to generate the scalar mesons $\sigma(600)$, $f_0(980)$, $a_0(980)$ and $\kappa(900)$ in the light sector, but with the $SU(4)$ symmetry strongly broken, mostly due to the explicit consideration of the masses of the vector mesons exchanged between pseudoscalars in the equivalent theory using the hidden gauge formalism for the vector mesons [46]-[49]. A different breaking of $SU(4)$ is also considered in [42], following general rules of $SU(n)$ breaking [50], which serves as an indication of theoretical uncertainties. We follow here the version based on the hidden gauge formalism, including the exchange of heavy (charmed) vector mesons in the Lagrangian.

We would like to put the $SU(4)$ breaking in a certain context. The basic assumption underlying [42, 45] is that the vertices in the hidden gauge formalism, of four vectors or three vectors, are approximately $SU(4)$ symmetric. As a first step, the kernel of the Bethe Salpeter equation (the potential) already breaks SU(4) symmetry in the terms that exchange heavy vectors, as we have mentioned. Later on, the amplitudes calculated with the unitary approach break $SU(4)$ symmetry because the physical masses of the particles are used to respect exactly thresholds and unitarity in coupled channels. This situation is already present in $SU(3)$. The starting lowest-order chiral Lagrangians are $SU(3)$ symmetric and the amplitudes obtained break $SU(3)$ symmetry due to the different masses of the particles belonging to the same $SU(3)$ multiplet. As an example, one of the $\Lambda(1405)$ states [51] and the $\Lambda(1670)$, which appear in the approach of [51, 52], are degenerate if the masses within the same $SU(3)$ multiplets are taken equal. One can see there an example of a large $SU(3)$ breaking present in nature, which, however, is consistent with assuming an exact $SU(3)$ meson-baryon interaction Lagrangian.

With these assumptions for the SU(4) symmetry and its breaking one gets realistic results for the spectra of mesons in [42, 45], which also agree with those obtained in

the heavy quark formalism for the case of light-heavy meson interaction [53], up to a mass term with no practical consequences. It is also interesting to note that the same basic $SU(4)$ assumptions are done for the interaction of mesons with baryons in [54, 55, 56] and one obtains realistic results concerning the $\Lambda_c(2595)$ and $\Sigma_c(2800)$ resonances.

Following [24], one derives the kernel (potential) from the lowest order Lagrangian and iterates it to generate all the terms of the Bethe Salpeter series, which can be summed up in the on shell formalism [57, 58] by means of

$$T = [1 - VG]^{-1}V \ , \tag{1}$$

where T is a matrix in the space of coupled channels representing the transition scattering amplitude from one channel to another and V the equivalent matrix for the transition potential. The diagonal matrix in the coupled channel space G accounts for the loop integral of the two particle propagator of any intermediate state

$$
\begin{aligned}
G_{ii} &= i \int \frac{dq^4}{(2\pi)^4} \frac{1}{q^2 - m_1^2 + i\epsilon} \frac{1}{(P-q)^2 - m_2^2 + i\epsilon} = \tag{2}\\
&\quad \frac{1}{16\pi^2} \left[\alpha_i + Log\frac{m_1^2}{\mu^2} + \frac{m_2^2 - m_1^2 + s}{2s} Log\frac{m_2^2}{m_1^2} \right.\\
&\quad + \left. \frac{\bar{p}}{\sqrt{s}} \left(Log\frac{s - m_2^2 + m_1^2 + 2\bar{p}\sqrt{s}}{-s + m_2^2 - m_1^2 + 2\bar{p}\sqrt{s}} + Log\frac{s + m_2^2 - m_1^2 + 2\bar{p}\sqrt{s}}{-s - m_2^2 + m_1^2 + 2\bar{p}\sqrt{s}} \right) \right] \ , \tag{3}
\end{aligned}
$$

where P is the total four momentum, q one of the pseudoscalar four momentum, $\bar{p}$ the on shell three momentum and m_1, m_2 the masses of the two pseudoscalars.

This integral requires dimensional regularization by means of a substraction constant, α_i, or cutting the three dimensional integral in G with a cut off. Both methods establish equivalent schemes in a certain chosen region of energies [58]. However, the use of the dimensional regularization method relies upon Lorentz covariance of magnitudes, which is lost in nuclei where one has a privileged reference frame, the one where the nucleus is at rest. As a result, the use of the dimensional regularization method introduces pathologies when including the selfenergy of the particles in the medium, which are avoided with the use of a cut off [59]. Hence, in the channels where we renormalize the particles in the medium, we stick to the cut off formalism. For all the other channels we use the same dimensional regularization approach of [42] which guarantees unitarity. The cut off method is used only in the channels DK and $D\bar{D}$, and the values of the cut off will be shown later on, but they are much bigger than the on-shell three momenta of the particles in the loops, such that the imaginary part of G_{ii} is obtained exactly in the free case. In free space we have

$$G_{ii}(s) = \int_0^{q_{max}} \frac{q^2 dq}{(2\pi)^2} \frac{\omega_1 + \omega_2}{\omega_1\omega_2[(P^0)^2 - (\omega_1 + \omega_2)^2 + i\epsilon]} \ , \tag{4}$$

where q_{max} stands for the cut off, $\omega_i = (\vec{q}_i^2 + m_i^2)^{1/2}$ and the center-of-mass energy $(P^0)^2 = s$. The expression of G_{ii} will be changed to account for the medium effects on the pseudoscalar mesons in Sec. 4.

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The T matrix of Eq. (1) generates poles for some quantum numbers. We look for them in the second Riemann sheet for the channels which are open and those poles are associated to resonances. For the channels where we found bound states, i.e., states below threshold, the poles appear in the first Riemann sheet. This is actually the case for $X(3700)$ in the $D\bar{D}$ and also for the $D_{0s}(2317)$ in the DK channel. Close to a pole, the amplitude looks like

$$T_{ij} \approx \frac{g_i g_j}{z - z_R} \ , \tag{5}$$

where Re z_R gives the mass of the resonance and Im z_R the half width. The constants g_i, obtained from the residues of the amplitudes, provide the coupling of the resonance to a particular channel and indicate the relevance of this channel in building up the resonance. This said, it is useful to recall that the $X(3700)$ and $D_{s0}(2317)$ are obtained from poles of the scattering matrix with quantum numbers ($C = 0$, $S = 0$, $I = 0$) and ($C = 1$, $S = 1$, $I = 0$), respectively. In Tables 1 and 2, we show the coupling of each state to the different channels contributing to these quantum numbers, which are obtained from [42] but including the η' and the mixing with η (see [60]).

Channel	Re(g_X) [MeV]	Im(g_X) [MeV]	$\lvert g_X \rvert$ [MeV]
$\pi^+\pi^-$	9	83	84
K^+K^-	5	22	22
D^+D^-	5962	1695	6198
$\pi^0\pi^0$	6	83	84
$K^0\overline{K^0}$	5	22	22
$\eta\eta$	1023	242	1051
$\eta\eta'$	1680	368	1720
$\eta'\eta'$	922	-417	1012
$D^0\overline{D^0}$	5962	1695	6198
$D_s^+D_s^-$	5901	-869	5965
$\eta_c\eta$	518	659	838
$\eta_c\eta'$	405	9	405

Table 1: $X(3700)$: Couplings of the pole at (3722-i18) MeV to the channels (C=0, S=0,I=0).

From these tables we observe the following:

1) The heavy singlet, hidden charm state $X(3700)$ couples most strongly to $D\bar{D}$. Next it couples to $D_s\bar{D}_s$. However, from the square of couplings which would enter into the selfenergy of the resonance due to the intermediate meson-meson states, one gets a factor of two more weight for the $D\bar{D}$ than the $D_s\bar{D}_s$ states. On the other hand, while the pole around 3700 MeV is close to the $D\bar{D}$ threshold of 3738 MeV, it is about 238 MeV away from the $D_s\bar{D}_s$ threshold of 3938 MeV. The off shellness of the $D_s\bar{D}_s$ in the loop function of Eq. (4) further reduces the contribution of this

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Channel	$\mathrm{Re}(g_{D_{s0}})$ [MeV]	$\mathrm{Im}(g_{D_{s0}})$ [MeV]	$|g_{D_{s0}}|$ [MeV]
K^+D^0	5102	0	5102
K^0D^+	5102	0	5102
ηD_s^+	-2952	0	2952
$\eta'D_s^+$	4110	0	4110
$\eta_C D_s^+$	2057	0	2057

Table 2: $D_{s0}(2317)$: Couplings of the pole at 2317 MeV to the channels (C=1, S=1,I=0).

channel to less than 10% of the $D\bar{D}$. This is important to note since, when dealing with the medium correction in the next sections, we shall consider the normalization of D but not the one of D_s.

2) The $D_{s0}(2317)$ couples most strongly to DK. The next channels are the $D_s\eta'$ and $D_s\eta$. The same considerations as before lead to a similar relative contribution of $D_s\eta'$ and $D_s\eta$ with respect to the dominant DK channel. Once again, we are lead to deal with the medium modification of the D meson, those of the K state being relatively unimportant since the KN interaction is not too strong and has no singularities [61].

It should be noted that in the step from $SU(3)$ to $SU(4)$ we introduce more uncertainties than one has in $SU(3)$. That is the reason why in [42] two different models, which break $SU(4)$ in different ways, were used to see the uncertainties in the results. It was found there that the results for the $X(3700)$ and $D_{s0}(2317)$ states were rather independent of the model, while other states predicted here were more model dependent. Even though, it is important to make an evaluation of the theoretical uncertainties for those states and, for this purpose, we have followed the same approach as in [42]. Thus, we evaluate the results using a sample of parameters which are varied within the reasonable limits discussed in [42]. In particular we vary the f_π and f_D parameters in the range $f_\pi \in [85, 115]$ MeV and $f_D \in [146, 218]$ MeV. This alone gives a good approximation to the uncertainties and we will perfom the medium calculations evaluating the uncertainties with this method.

The main results of the paper are the effects of the medium in the $D_{s0}(2317)$. The results for this resonance were obtained in [42], where the prescription of using f_π for the light mesons and f_D for the heavy ones was used. For the case of scattering of light mesons with heavy ones, chiral symmetry requires the use of f_π in all cases (eventually f_K if kaons are involved). For this reason we redo the calculations with the new chiral symmetric prescription. We also introduce the novelty with respect to [42] of considering also the η' in the set of pseudoscalar mesons according to the method of [60]. We should bear in mind that the substraction constant of Eq. (3) was slightly tuned from its natural value to get the mass of the $D_{s0}(2317)$ at the right place. Hence, we do the same here and then we look at the results obtained for the couplings to the channels in the two cases. Those are shown in Table 3. As we can see, the results are very similar. For the most important building block, the DK channel, the differences of the coupling are of 4%, indicating that the errors induced

Channel	$\lvert g_i \rvert$ Model A [MeV]	$\lvert g_i \rvert$ Model B [MeV]
DK	7215	7503
$D_s\eta$	2952	3005
$D_s\eta\prime$	4110	4146
$D_s\eta_c$	2058	1246

Table 3: Couplings of the $D_{s0}(2317)$ to its building blocks. Model A refers to the model using both, f_π and f_D in the couplings, while in Model B only f_π is used, respecting constrains from chiral symmetry. The channels are in isospin basis. The position of the pole is fixed in both models to 2317 MeV, taking $\alpha_H = -1.48$ in the model A, and $\alpha_H = -1.16$ in the model B (α_H means the substraction constant used in [42] for the channels involving at least one heavy pseudoscalar meson).

by the explicit chiral symmetry breaking of [42] are very small. Yet, in the present paper we shall use the chiral symmetric version described here. Using f_K instead of f_π leads to much smaller differences than those in Table 3. However, we shall use just f_π.

For the $X(3700)$, which comes mostly from $D\bar{D}$, we have no such constraints from chiral symmetry and we follow the approach of [42], except for the inclusion of the $\eta\prime$ channel. We evaluate uncertainties in the results, though, by using again f_π instead of both f_D and f_π. The stability of the couplings with respect to the changes done can be partly justified by recalling that, with one channel dominance, the coupling of a bound state to its constituents is given in terms of the binding energy by the compositness condition of Weinberg [62]-[67]. In the present case, the DK is the dominant channel, but other channels also matter. This is why some changes, although small, were found in the couplings while demanding the same binding energy.

3 The selfenergy of the D meson

3.1 s-wave selfenergy

We shall use the $T = 0$ results from the work of [56]. There, the D meson selfenergy is obtained from a selfconsistent coupled channel calculation, whose driving term is a broken SU(4) s-wave Weinberg Tomozawa interaction supplemented by an attractive isoscalar-scalar term. The introduction of a supplementary scalar-isoscalar interaction in the diagonal DN channel, the Σ_{DN} term, which is prevalent in the QCD sum rule and mean-field approaches, is, however, a subject of controversy. Its effect on the D meson self-energy was studied in [55] and [56], and compared to the case where this term was neglected. It was found in [56] that the results obtained with or without this term of uncertain origin were qualitatively identical, such that the phenomenology did not allow to draw any conclusion concerning it. The differences found there between the two options are far smaller than the uncertainties that we

have in the present problem from other sources and, hence, we take the option of ignoring this term. The Bethe Salpeter equation is then solved using a cut off regularization, which is fixed by reproducing the position and width of the $\Lambda_c(2593)$ resonance. As a result, a new resonance in the $I = 1$ channel is generated around 2800 MeV, the $\Sigma_c(2800)$. The coupled channel structure includes: $\pi\Lambda_c$, $\pi\Sigma_c$, DN, $\eta\Lambda_c$, $K\Xi_c$, $\eta\Sigma_c$, $K\Xi_c'$, $D_s\Lambda$, $D_s\Sigma$, $\eta'\Lambda_c$ and $\eta'\Sigma_c$.

The in medium s-wave DN amplitude accounts for Pauli blocking effects on the nucleons in the DN channel, mean-field bindings of baryons via a σ-ω model, and renormalization of π and the D through their corresponding selfenergies in the intermediate propagators. The s-wave D selfenergy is obtained iteratively following a selfconsistent procedure as one integrates the in medium s-wave DN amplitude over the nucleon Fermi sea $n(p)$:

$$\Pi_D^{(s)}(q^0, \vec{q}, \rho) = \int \frac{d^3p}{(2\pi)^3}\, n(p)\, [\tilde{T}_{DN}^{(I=0)}(P^0, \vec{P}, \rho) + 3\tilde{T}_{DN}^{(I=1)}(P^0, \vec{P}, \rho)]\;, \tag{6}$$

where $\tilde{T}^{(I=0,1)}$ stands for the in medium s-wave DN amplitudes in $I = 0$ and $I = 1$. The quantities $P^0 = q^0 + E_N(\vec{p})$ and $\vec{P} = \vec{q} + \vec{p}$ are the total energy and momentum of the DN pair in the nuclear matter rest frame, with $E_N(\vec{p})$ being the single-particle nucleon energy and the values $(q^0, \vec{q})$ the energy and momentum of the D meson also in this frame. For more details see [56].

In fact, the s-wave D meson selfenergy in the medium was initially studied in [68, 69, 54] and further work was done in [55, 56]. There are novelties in the approach of [55, 56] with respect to the earlier works. Indeed, for consistency with the reduction from t-channel vector-meson exchanges to a zero-range Weinberg-Tomozawa form, corroborated with explicit cancellations of terms, the model of [55, 56] removes a factor $k^\mu k^\nu / M_V^2$, which was used in [54], and by means of which a better width for the $\Lambda_c(2593)$ is obtained in [55, 56]. Also in [55, 56] the authors use a conventional momentum cut-off regularization that was found to be more appropiate than the dimensional method in view of its application to meson-baryon scattering in the nuclear medium where Lorentz covariance is manifestly broken, as discussed in Section 2.

With respect to the work of [68, 69] there are also novelties in [55, 56]. In the exploratory work of [68, 69], the free amplitudes were constructed from separable coupled channel interactions obtained from chiral motivated Lagrangians upon replacing the s quark by the c quark. While these works give the first indication that the $\Lambda_c(2593)$ could have a dynamical origin, they ignored the strangeness degree of freedom due its very construction. Therefore, the π and K (Goldstone) mesons were not treated on an equal footing, and the role of some channels that would appear in the corresponding SU(4) meson and baryon multiplets was ignored.

3.2 p-wave selfenergy

We start by recalling the $SU(3)$ chiral Lagrangian [70, 71] for the coupling of pseudoscalar mesons of the octet of the π to the baryon octet of the proton p

$$\mathcal{L}_1^{(B)} = \frac{1}{2}D < \bar{B}\gamma^\mu\gamma_5\{u_\mu, B\} > + \frac{1}{2}F < \bar{B}\gamma^\mu\gamma_5[u_\mu, B] >\;, \tag{7}$$

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where

$$u^2 = U = e^{i\frac{\sqrt{2}\phi}{f}} , \tag{8}$$

with ϕ the usual $SU(3)$ matrix of the meson fields, $f = 1.15 f_\pi$ with $f_\pi = 93$ MeV and

$$u_\mu = iu^\dagger \partial_\mu U u^\dagger = -\frac{\sqrt{2}}{f}\partial_\mu\phi + O(\phi^3) . \tag{9}$$

The B and $\bar{B}$ terms stand for the $SU(3)$ matrices of the baryon fields and $<>$ for the trace in $SU(3)$. Hence, at the one meson field level we have

$$\mathcal{L}_1^{(B)} = -\frac{1}{\sqrt{2}f}D < \bar{B}\gamma^\mu\gamma_5\{\partial_\mu\phi, B\} > -\frac{1}{\sqrt{2}f}F < \bar{B}\gamma^\mu\gamma_5[\partial_\mu\phi, B] > . \tag{10}$$

The $\bar{u}(\vec{p}')\gamma^\mu\gamma_5 u(\vec{p})$ vertex, assuming $\vec{p} \simeq 0$ since it will be the momentum of a nucleon in the Fermi sea, can be expressed up to $O(1/M^2)$ in terms of the $\vec{\sigma}$ operator such that the T matrix corresponding to the diagram of Fig. 1 is given by
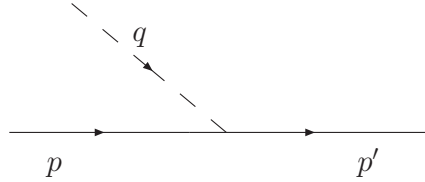
Figure 1: Meson-baryon scattering with an outgoing baryon. The labels p, p' and q refer to the momenta of the initial baryon, final baryon and meson respectively.

$$-it = \frac{1}{\sqrt{2}f}\vec{\sigma}\cdot\vec{q}\left(1 - \frac{q^0}{2M'}\right)[(D+F) < \bar{B}\phi B > +(D-F) < \bar{B}B\phi >] , \tag{11}$$

with M' the mass of the outgoing baryon in Fig. 1. We take $D = 0.80$ and $F = 0.46$ from [72, 73, 74]. In order to evaluate the coupling of the D meson to the nucleon and Λ_c, Σ_c we use $SU(4)$ symmetry. We couple the 20-plet of the baryons, to which the nucleon belongs, to the $\overline{20}$ representation of the antibaryons in order to give the 15-plet of the mesons of the π and the D [75]. By using the $SU(4)$ Clebsch-Gordan coefficients of [76], we have two independent irreducible matrix elements which can be related to the D and F coefficients. The result is that the couplings $D^0 p \to \Lambda_c^+$, $D^0 p \to \Sigma_c^+$, $D^+ p \to \Sigma_c^{++}$, $D^0 n \to \Sigma_c^0$, $D^+ n \to \Lambda_c^+$, $D^+ n \to \Sigma_c^+$ are identical to those of $K^- p \to \Lambda$, $K^- p \to \Sigma^0$, $\bar{K}^0 p \to \Sigma^+$, $K^- n \to \Sigma^-$, $\bar{K}^0 n \to \Lambda$, $\bar{K}^0 n \to \Sigma^0$ given in [77] by

$$-iV_{DNY} = \vec{\sigma}\cdot\vec{q}\left(1 - \frac{q^0}{2M'}\right)\left[\alpha\frac{D+F}{2f} + \beta\frac{D-F}{2f}\right] , \tag{12}$$

with the coefficients α, β of the Table 4.

We also take into account the coupling of the D meson with $\Sigma_c^*(2520)$ and N, in analogy to the p-wave interaction of pions and kaons with nucleons. For pions and

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	$D^0 p \to \Lambda_c^+$	$D^0 p \to \Sigma_c^+$	$D^0 n \to \Sigma_c^0$	$D^+ n \to \Lambda_c^+$	$D^+ n \to \Sigma_c^+$	$D^+ p \to \Sigma_c^{++}$
α	$-\frac{2}{\sqrt{3}}$	0	0	$-\frac{2}{\sqrt{3}}$	0	0
β	$\frac{1}{\sqrt{3}}$	1	$\sqrt{2}$	$\frac{1}{\sqrt{3}}$	-1	$\sqrt{2}$

Table 4: Coefficients for the DNY couplings

kaons it was shown that the $N^{-1}\Delta$ and $N^{-1}\Sigma^*(1385)$ excitations, respectively, are relevant for the calculation of the p-wave self-energy. Once again, we obtain the same result as in [77] for the $N^{-1}\Sigma^*(1385)$

$$- iV_{DNY^*} = a\vec{S}^\dagger \cdot \vec{q} \left(\frac{2\sqrt{6}}{5} \frac{D+F}{2f} \right) , \tag{13}$$

with $\vec{S}^\dagger$ being the spin $1/2 \to 3/2$ transition operator and a the coefficients given in Table 5.

	$D^0 p \to \Sigma_c^{*+}$	$D^0 n \to \Sigma_c^{*0}$	$D^+ p \to \Sigma_c^{*++}$	$D^+ n \to \Sigma_c^{*+}$
a	$-\frac{1}{\sqrt{2}}$	-1	-1	$\frac{1}{\sqrt{2}}$

Table 5: Coefficient for the $DN\Sigma_c^*(2520)$ couplings
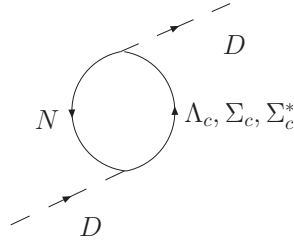
Figure 2: p-wave selfenergy diagram of the D meson.

With all those couplings, we can readily evaluate the p-wave D selfenergy given by the diagram of the Fig. 2, in complete analogy to [77]. The p-wave contribution coming from the $N^{-1}\Lambda_c$ and $N^{-1}\Sigma_c$ excitations reads

$$\Pi_{D^0}^{(p)}(q^0, \vec{q}, \rho) = \{\frac{1}{2} B_{D^0 p \Lambda_c^+}^2 \vec{q}^2 U_{\Lambda_c^+}(q^0, \vec{q}, \rho) + \frac{1}{2} B_{D^0 p \Sigma_c^+}^2 \vec{q}^2 U_{\Sigma_c^+}(q^0, \vec{q}, \rho)$$
$$+ \frac{1}{2} B_{D^0 n \Sigma_c^0}^2 \vec{q}^2 U_{\Sigma_c^0}(q^0, \vec{q}, \rho)\} F_L^2(q^2) , \tag{14}$$

where

$$B_{DNY} = \left(1 - \frac{q^0}{2M_Y}\right) \left[\alpha \frac{D+F}{2f} + \beta \frac{D-F}{2f}\right] , \tag{15}$$

10

and U is the Lindhard function for the $N^{-1}Y$ excitation given by

$$
\begin{aligned}
\operatorname{Re} U_Y(q^0, \vec{q}, \rho) &= \frac{3}{2}\rho\frac{M_Y}{qp_F}\left\{z + \frac{1}{2}(1-z^2)\ln\frac{|z+1|}{|z-1|}\right\} \\
\operatorname{Im} U_Y(q^0, \vec{q}, \rho) &= -\pi\frac{3}{4}\rho\frac{M_Y}{qp_F}\left\{(1-z^2)\theta(1-|z|)\right\} \\
z &= \left(q^0 - \frac{q^2}{2M_Y} - (M_Y - M)\right)\frac{M_Y}{qp_F} ,
\end{aligned}
\tag{16}
$$

with $\rho = \rho_n + \rho_p$, the nuclear density, $p_F = (3\pi^2\rho/2)^{1/3}$ the Fermi momentum, M_Y the hyperon mass and M the nucleon mass. The same result holds for the p-wave D^+ selfenergy ignoring small mass differences between particles of the same isospin multiplet.

The p-wave selfenergy due to the excitation of the decuplet is also readily evaluated and we find

$$
\Pi_{D^0}^{(p)*}(q^0, \vec{q}, \rho) = \left\{\frac{1}{3}C^2_{D^0 p\Sigma_c^{+*}}\vec{q}\,^2 U_{\Sigma_c^{+*}}(q^0, \vec{q}, \rho) + \frac{1}{3}C^2_{D^0 n\Sigma_c^{0*}}\vec{q}\,^2 U_{\Sigma_c^{0*}}(q^0, \vec{q}, \rho)\right\}F_L^2(q^2) ,
\tag{17}
$$

where

$$
C_{DNY} = af_Y^*\frac{2\sqrt{6}}{5}\frac{D+F}{2f} ,
\tag{18}
$$

with a given in Table 5 and f_Y^* being a recoil factor [77], which we approximate by $f_Y^* \simeq (1 - M_D/M_Y)$.

In Eqs. (14) and (17), we include a form factor of monopole type at the D meson-baryon vertices by analogy to the one accompanying the Yukawa πNN vertex [77]-[80].

$$
F_L(q^2) = \frac{\Lambda^2}{\Lambda^2 + \vec{q}\,^2} \quad \text{with } \Lambda = 1.05 \text{ GeV.}
\tag{19}
$$

This form factor is suited for light hadrons, i.e., pion excitation of ph. However, it is unlikely that the range of Λ is the same when dealing with D mesons. There are indications that the form factor to account for off shell D mesons requires a value of Λ substantially larger [81]. We shall come back to this point at the end of the Results Section, reevaluating results with the heavy meson form factor and analyzing the uncertainties.

With regards to the p-wave D^+ selfenergy, it turns out to be the same as for D^0 in symmetric nuclear matter $\rho_n = \rho_p$.

For the $\bar{D}$ meson, we note that the p-wave $\bar{D}$ selfenergy would correspond to the diagrams in Fig. 3, which involve the difference between the sum of $\bar{D}$ and Y masses, and the nucleon mass. The contribution of those diagrams is negligible due to the large mass of the $\bar{D}$ and Σ_c, Λ_c. The same holds for the p-wave $\bar{D}$ selfenergy coming from the $N^{-1}\Sigma_c^*(2520)$ excitation.

Figure 3: p-wave selfenergy diagrams of the $\bar{D}$ meson.

3.3 The D meson spectral function

The selfenergy of a D meson in nuclear matter is given by the coherent sum of the s-wave and p-wave selfenergies:

$$\Pi_D(q^0, \vec{q}, \rho) = \Pi_D^{(s)}(q^0, \vec{q}, \rho) + \Pi_D^{(p)}(q^0, \vec{q}, \rho) + \Pi_D^{(p)*}(q^0, \vec{q}, \rho) . \tag{20}$$

Then, the D propagator is written in the medium as

$$D_D(q^0, \vec{q}, \rho) = \frac{1}{(q^0)^2 - \vec{q}^2 - m_D^2 - \Pi_D(q^0, \vec{q}, \rho)} . \tag{21}$$

For later purposes, it is convenient to write the D propagator in terms of its Lehmann representation [82, 83]

$$D_D(q^0, \vec{q}, \rho) = \int_0^\infty d\omega \left\{ \frac{S_D(\omega, \vec{q}, \rho)}{q^0 - \omega + i\eta} - \frac{S_{\bar{D}}(\omega, \vec{q}, \rho)}{q^0 + \omega - i\eta} \right\} , \tag{22}$$

where S_D and $S_{\bar{D}}$ are the spectral functions of D and $\bar{D}$, respectively,

$$S_{D(\bar{D})}(q^0, \vec{q}, \rho) = -\frac{1}{\pi} \frac{\text{Im}\Pi_{D(\bar{D})}(q^0, \vec{q}, \rho)}{\left| (q^0)^2 - \vec{q}^2 - m_D^2 - \Pi_{D(\bar{D})}(q^0, \vec{q}, \rho) \right|^2} . \tag{23}$$

In the calculations we will ignore the selfenergy of the $\bar{D}$. As in the case of the K with respect to the $\bar{K}$, the $\bar{D}$ selfenergy is smaller than that of the D meson and, more importantly, it has no imaginary part from inelastic channels. Hence, it does not lead to modifications of the width of the states that we study here, which is the most striking change that we find. We will discuss this in more detail in what follows.
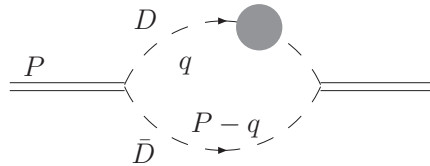

Figure 4: (Color online) The $D\bar{D}$ loop function of the scalar meson. The shaded circle indicates the D selfenergy insertion.
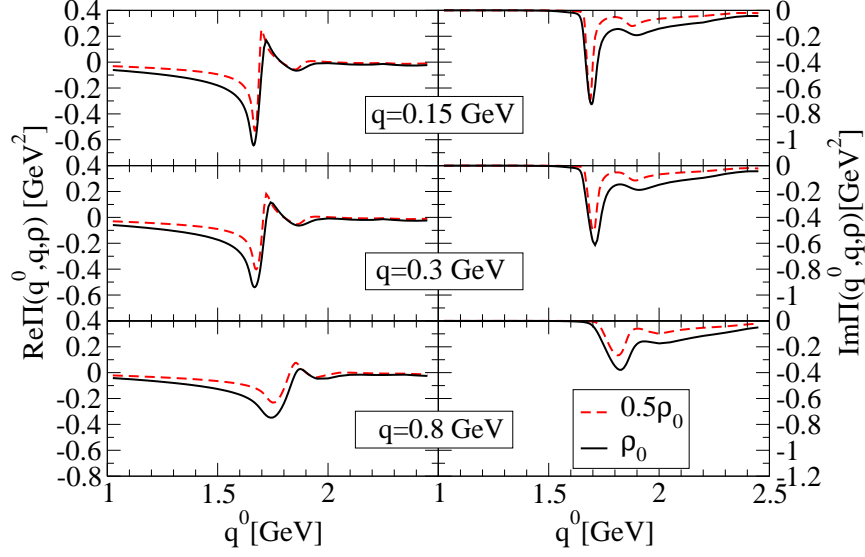
Figure 5: (Color online) Real (left column) and Imaginary (right column) part of the s-wave D selfenergy for D three-momenta $q = 0.15, 0.3$ and 0.8 GeV as a function of the D energy q^0 at densities $\rho = 0.5\rho_0$ and $\rho = \rho_0$, with $\rho_0 = 0.17\text{fm}^{-3}$ the normal nuclear matter density.
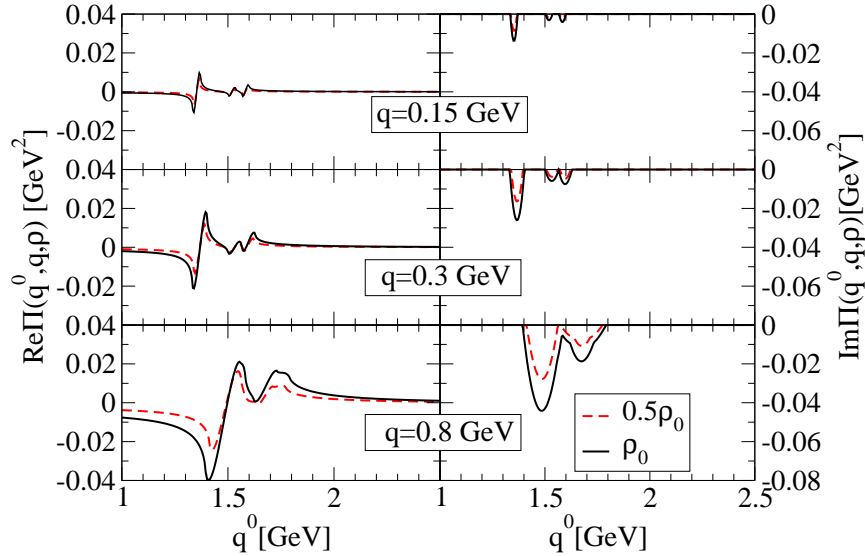

Figure 6: (Color online) Real (left column) and Imaginary (right column) part of the p-wave D selfenergy for D three-momenta $q = 0.15, 0.3$ and 0.8 GeV as a function of the D energy q^0 at densities $\rho = 0.5\rho_0$ and $\rho = \rho_0$, with $\rho_0 = 0.17\text{fm}^{-3}$ the normal nuclear matter density.

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4 Two meson loop function in the medium

The medium modifications are introduced in the two loop meson function by using the dressed two meson propagator in nuclear matter. As an example, let us evaluate the G function in the medium for the $D\bar{D}$ intermediate state:

$$\widetilde{G}(P^0, \vec{P}, \rho) = i \int \frac{d^4q}{(2\pi)^4} D_D(q, \rho) D_{\bar{D}}(P - q, \rho) \ . \tag{24}$$

We shall dress the D propagator and leave the $\bar{D}$ propagator free. The reason to neglect the $\bar{D}$ selfenergy in the medium is that it is very small compared to its mass [56]. Indeed, the p-wave part is negligible as discussed at the end of Subsection 4.2. The s-wave part is equally small, but more importantly there is no absorption of $\bar{D}$ by nucleons, meaning that the $\bar{D}N$ does not decay to baryonic resonances which have c quarks instead of $\bar{c}$. The analogy is clear with the K and $\bar{K}$, where the $\bar{K}$ (analogous to D) can undergo absortion reactions $\bar{K}N \to \pi\Lambda$, $\pi\Sigma$, while the K^+ (analogous to $\bar{D}$) cannot be absorbed. Altogether this justifies to use the free propagator for $\bar{D}$. We must evaluate the loop function of the diagram of Fig. 4, where the blob in the D propagator symbolizes the D selfenergy insertion, indicating that we must use the D propagator in the medium. Then, we use Eq. (22) for this propagator. Given the large mass of the D mesons, we can also neglect the $S_{\bar{D}}$ part in the propagator of Eq. (22), which upon q^0 integration in the Eq. (24) leads to a contribution of order of $1/(P^0 + 2\omega_D(q))$, very small compared with the part coming from S_D. Thus, we can write

$$
\begin{aligned}
\widetilde{G}(P^0, \vec{P}, \rho) &= i \int \frac{d^4q}{(2\pi)^4} \int_0^\infty d\omega \frac{S_D(\omega, \vec{q}, \rho)}{q^0 - \omega + i\eta} \frac{1}{(P^0 - q^0)^2 - \vec{q}^2 - m_{\bar{D}}^2 + i\eta} \\
&= \int \frac{d^3q}{(2\pi)^3} \int_0^\infty d\omega \frac{S_D(\omega, \vec{q}, \rho)}{P^0 - \omega - \omega_{\bar{D}}(\vec{q}) + i\eta} \frac{1}{2\omega_{\bar{D}}(\vec{q})} \ ,
\end{aligned}
\tag{25}
$$

where $\omega_{\bar{D}}(\vec{q}) = (\vec{q}^2 + m_{\bar{D}}^2)^{1/2}$. We evaluate Eq. (25) with a three momentum cut off of $q_{max} = 0.85$ and 0.9 GeV for the $X(3700)$ and $D_{s0}(2317)$, respectively, equivalent to the use of dimensional regularization of [42] with the chosen substraction constants.

Following the discussion in the former section, the $D\bar{D}$ loop function will appear in the case of the $X(3700)$ hidden charm state. In the case of the $D_{s0}(2317)$, we have K instead of $\bar{D}$ and, again, we use its free propagator neglecting the small structureless K selfenergy.

5 Results

In Figs. 5 and 6 we show the results of the s-wave and p-wave selfenergies of the D meson. We perform the calculations first by using the light meson version of the Dph form factor, Eq. (19). We will show results with a heavy meson form factor at the end of this section. In Fig. 5 the structures around energies 1.7 GeV and 2 GeV correspond to the excitation of the $h\Lambda_c(2593)$ and $h\Sigma_c(2800)$, where the $\Lambda_c(2593)$ and $\Sigma_c(2800)$ are $1/2^{*-}$ dynamically generated resonances of the theory and h stands for

hole of nucleon. With regard to the p-wave contribution, the structures seen in the p-wave selfenergy, with peaks for the imaginary part and oscillations in the real one around 1.4 to 1.8 GeV, correspond to the excitation of the different hY components, with $Y = \Lambda_c, \Sigma_c, \Sigma_c^*$.

Note that the p-wave selfenergy is much smaller than the s-wave selfenergy, even up to momentum as large as 800 MeV/c. The main reason is that the s-wave selfenergy of the mesons goes roughly as the meson mass, from the Weinberg Tomozawa interaction, while the p-wave scales differently, roughly like the baryon mass. In the case of pions, the p-wave selfenergy was more important than the s-wave [84], but in the case of kaons the relative importance of the p-wave was already smaller [59, 85].

In Figs. 7 and 8, we plot the $\widetilde{G}(P^0, \vec{P}, \rho)$ function for $\vec{P} = 0$ for the $D\bar{D}$ and DK loops for different densities. The effects of the density in the loop function are clearly visible in all cases. The imaginary part is largely increased at lower energies and collects strength below threshold of the $D\bar{D}$ or DK channels, respectively, due to the opening of new many body decay channels of the meson-meson system. As an example, let us take the DK loop function (see Fig. 9). The D is renormalized and accounts for $DN \to \pi\Lambda_c$, $\pi\Sigma_c$ or $DN \to \Lambda_c$, Σ_c. Hence, the DK loop in the medium accounts for intermediate channels $h\pi\Lambda_c K$, $h\pi\Sigma_c K$ or $h\Lambda_c K$, $h\Sigma_c K$ which have a smaller mass than the DK system and open up at lower energies than DK threshold. The real parts of $\widetilde{G}$ are also sizeably modified around the thresholds as one can see in the figures.
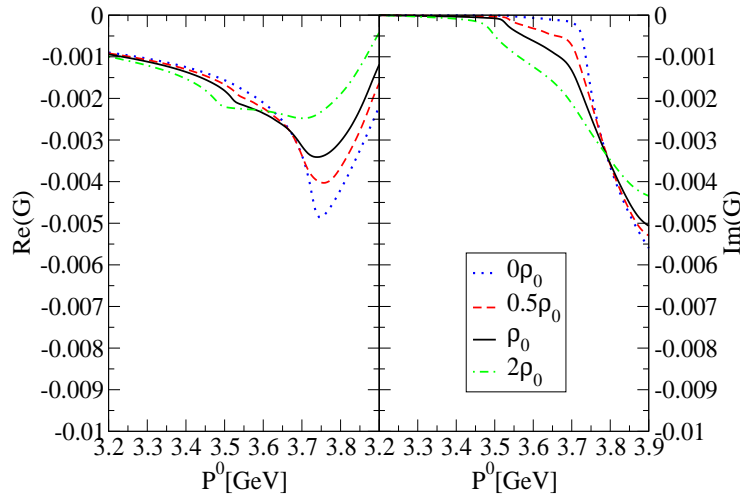

Figure 7: (Color online) Loop function in the medium: $\mathrm{Re}\,\widetilde{G}(P^0, \vec{P}, \rho)$ (left) and $\mathrm{Im}\,\widetilde{G}(P^0, \vec{P}, \rho)$ (right) for $D\bar{D}$, the channel with the largest coupling to the X(3700) meson. $\widetilde{G}(P^0, \vec{P}, \rho)$ is given from Eq. (25).

In Fig. 10 we show the $|T|^2$ for the $D^0 K^+ \to D^0 K^+$ amplitude around the region of 2300 MeV for different densities. We can see that originally, at $\rho = 0$, the amplitude exhibits the pole of the $D_{s0}(2317)$, with zero width. This is the reason why it is out of the y-scale in the plot since $|T|^2$ goes up to infinite. As the density increases, we can see a slight shift of the mass, of the order of 15 MeV attraction at $\rho = \rho_0$.
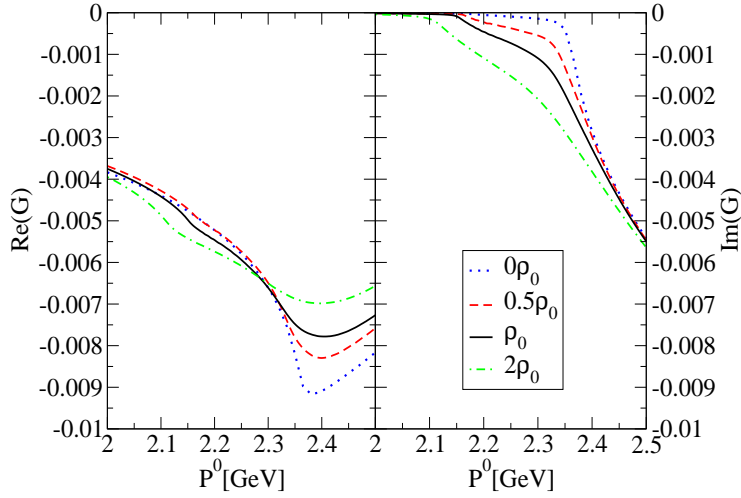
Figure 8: (Color online) Loop function in the medium: $\text{Re}\,\widetilde{G}(P^0,\vec{P},\rho)$ (left) and $\text{Im}\,\widetilde{G}(P^0,\vec{P},\rho)$ (right) for DK, the channel with the largest coupling to the $D_{s0}(2317)$ meson. $\widetilde{G}(P^0,\vec{P},\rho)$ is given from Eq. (25).

However, the increase in the width is more spectacular, which goes from zero in the free case to about 100 MeV at $\rho = \rho_0$, and 200 MeV at $\rho = 2\rho_0$. This is certainly a drastic relative effect, and even big in absolute value. The origin of the width in the medium is due to the opening of new channels $DN \to \Lambda_c, \Sigma_c$ from the p-wave selfenergy and $DN \to \pi\Lambda_c, \pi\Sigma_c$ from the s-wave selfenergy. On the other hand, the use of selfconsistency in the evaluation of the D selfenergy [55, 56] generates also some two nucleon induced D absorption channels like $DNN \to N\Lambda_c, \pi N\Lambda_c, \pi N\Sigma_c$, etc. All these channels collaborate to make the D disappear inside the nuclear medium through DN or DNN inelastic reactions, where the D gets absorbed.
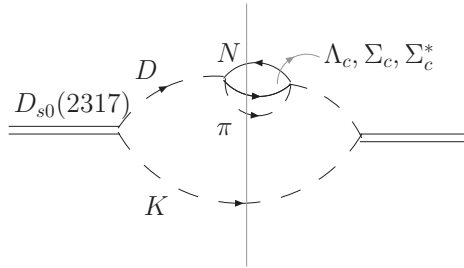

Figure 9: Decay channel of the $D_{s0}(2317)$ in the nucleus into $K\pi\Lambda_c$ or $K\pi\Sigma_c$.

The study of such decay channels in the nucleus would offer information on the coupling of the $D_{s0}(2317)$ to the DK, the basic building block of the resonance according to the underlying theory that we are using. As an example, the exploration of the decay channel of the $D_{s0}(2317)$ in the nucleus into $K\pi\Lambda_c$ or $K\pi\Sigma_c$ channels,
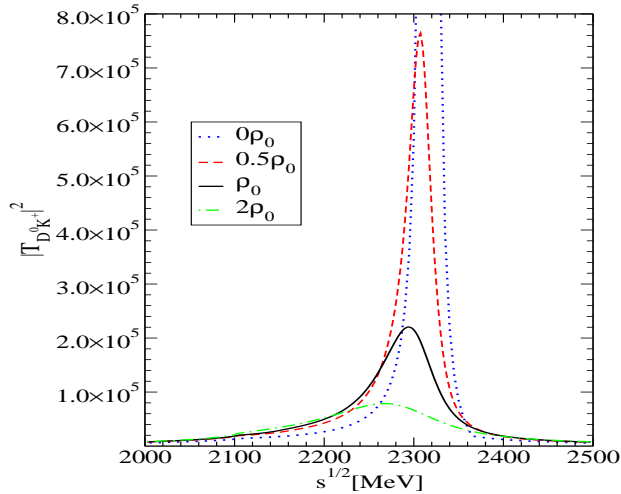
16

Figure 10: (Color online) $D_{s0}(2317)$ resonance: $|T|^2$ for the $D^0K^+ \to D^0K^+$ amplitude for different densities.

which corresponds to the cut in the diagram of Fig. 9, would provide combined information on the $D_{s0}(2317) \to DK$ coupling and the $DN \to \pi\Lambda_c(\Sigma_c)$ cross section.

The results with the large width for the $D_{s0}(2317)$ are a consequence of the large coupling of this resonance to DK. This large coupling is guaranteed by the "compositness condition" of Weinberg [62]-[67], as far as the resonance is dynamically generated and DK is the main building block. Should this resonance be a $q\bar{q}$ state as suggested in [86] or have any other structure, like $q\bar{q}$ with a mixture of DK components as suggested in [87, 88], such large coupling would not appear [62] and, thus, the width in the medium would be much smaller than we predict here. Hence, investigating the widths in the medium provides extra information on the nature of these resonances.

In Fig. 11 we show the same results for the hidden charm $X(3700)$ resonance. The resonance begins with about 60 MeV in the real axis (36 MeV deduced from the imaginary part of the pole position) from its main decay into the $\eta\eta$, $\eta\eta'$, and $\eta'\eta'$ channels. The difference with [42], where the width was much smaller, is due to the fact that, here, we use the mixing of mesons η and η' [60], which makes the η_c a pure $c\bar{c}$ state rather than pure $SU(4)$ state of the 15-plet of mesons. The study of this figure shows that at $\rho = \rho_0$ the width has become as large as 250 MeV and, due to cusps on the multiparticle channels that open up, the strength of the resonance in the nuclear medium acquires a peculiar shape as the density increases. Here, the main decay channels in the nucleus are the same as in the case of the $D_{s0}(2317)$, replacing the K by a $\bar{D}$. One should note that there is no relationship of the important decay channels to those that dominate the dynamical generation of the state. For example, the $X(3700)$ couples mainly to $D\bar{D}$ but the $D\bar{D}$ threshold is above the $X(3700)$ in free space and there is no decay into these components. The analogy can be made with the $f_0(980)$ which couples mostly to $K\bar{K}$ [89, 90] but decays into $\pi\pi$ in free space, since there is no phase space for $K\bar{K}$ decay. The coupling of the $X(3700)$
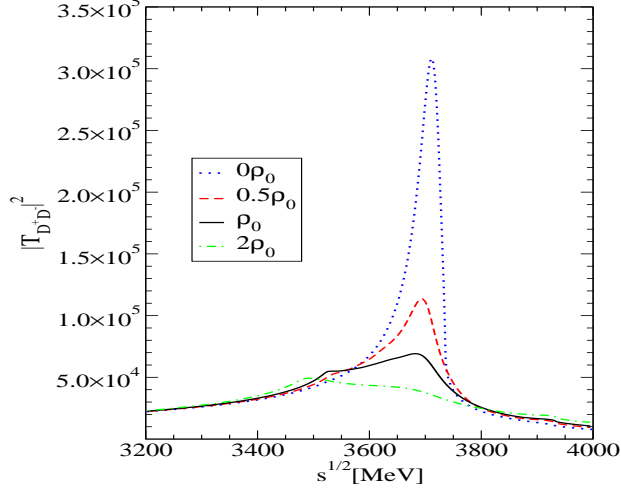
17

Figure 11: (Color online) $X(3700)$ resonance: $|T|^2$ for the $D^0\bar{D}^0 \rightarrow D^0\bar{D}^0$ amplitude for different densities.

to the $\eta\eta$, $\eta\eta'$ etc components is small, reflecting the fact that these channels are relatively unimportant in the structure of the resonance. However, because there is large phase space for these decays, they are mostly responsible for the $X(3700)$ width.

In the medium things might be different. Indeed, in the presence of nucleons we can have $DN \rightarrow \pi\Lambda_c$, $\pi\Sigma_c$ as we mentioned above, and new decay channels for the $X(3700)$ appear, such as $X(3700)N \rightarrow \bar{D}\pi\Lambda_c$, $\bar{D}\pi\Sigma_c$ [see Fig. 9 (changing K by $\bar{D}$)], which have nearly 400 MeV phase space available. The strong coupling of the $X(3700)$ to $D\bar{D}$ and the opening of these new decay modes makes now the width in the medium sizeable. We note that the trend of changes of $|T|^2$ with nuclear density is typically observed in resonance properties with temperature and/or density [91, 92, 93].

As mentioned at the end of Subsection 4.2, in order to analyze the uncertainties linked to the use of different form factors, we reevaluate the results using a form factor for the off shell D mesons which we obtain from the work of [81]. The heavy meson form factor is now

$$F_H(q^2) = \frac{\Lambda_D^2 - m_D^2}{\Lambda_D^2 - q^2} \quad \text{with } \Lambda_D = 3.5\,\text{GeV}. \tag{26}$$

The value of q^2 in $F_H(q^2)$ is taken for the configurations that give rise to the width in diagram of Fig. 9 when the D selfenergy comes from particle-hole (ph) excitation of Fig. 2. This means that, for the case of the $D_{s0}(2317)$, we place the K and $\Lambda_c h$, $\Sigma_c h$ or $\Sigma_c^* h$ on shell. This occurs for any value of the running variable q at a value of q^0 given by

$$q^0 = M_{D_{s0}} - \omega_K(q_{on}) \quad \text{with } q_{on} = \frac{\lambda^{1/2}((M_{D_{s0}} + m_N)^2, M^2, m_K^2)}{2\,(M_{D_{s0}} + m_N)} \tag{27}$$

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ρ	M[MeV]	Γ[MeV]
0.0	2316±5	0
0.5	2306±17	58±10
1.0	2295±23	115±25
1.5	2283±25	150±25
2.0	2274±31	190±30

Table 6: Mass and width for the $D_{s0}(2317)$ at different densities with error bands due to the uncertainties of our model.

ρ	M[MeV]	Γ[MeV]
0.0	3710±18	60± 10
0.5	3691± 10	135±20
1.0	3638± 15	255±25
1.5	3599±15	320±25
2.0	3565±29	340±25

Table 7: Mass and width for the $X(3700)$ at different densities with error bands due to the uncertainties of our model.

where M can be M_{Λ_c}, M_{Σ_c} or $M_{\Sigma_c^*}$ depending on the type of ph excitation. We anticipate that, given the minor relevance of the p-wave selfenergy, the differences with respect to the former calculation using the light meson form factor will be small. This is indeed the case as can be seen in Fig. 12, where we show $|T|^2$ for the case of the $D_{s0}(2317)$ for the two types of form factor. As we can see there are small differences in the position of the peak and no difference in the width.

As indicated before, in Figs. 10 and 11 we have shown the squared amplitudes for different densities to facilitate the comparison. As indicated at the end of Section 2, we also evaluate the uncertainties in the results. We show in Tables 6 and 7 the results obtained for the mass and width of the resonances at different densities, with their uncertainties obtained from a Monte Carlo sampling of the parameters f_π and f_D as indicated in Section 2 and evaluated from the plots of $|T|^2$ in the real axis. The results in these tables are all evaluated using the heavy meson form factor of Eq. (26). The relevant information from these tables is that the differences between the widths at $\rho = 0$ or $\rho = \rho_0$ for both resonances are much bigger than the uncertainties in the width from uncertainties in the model. In the case of the masses, we do not see appreciable shift in the mass compared to the uncertainties for the case of the $D_{s0}(2317)$, see Table 6. For the $X(3700)$, the shift from $\rho = 0$ to $\rho = \rho_0$ is of the order of 70 MeV, see Table 7, smaller than the width in the medium. Hence we cannot make any strong point concerning a possible shift of the masses.

In view of future experiments measuring medium modifications of these resonances, we can recall the method that has proved most efficient in measuring nuclear
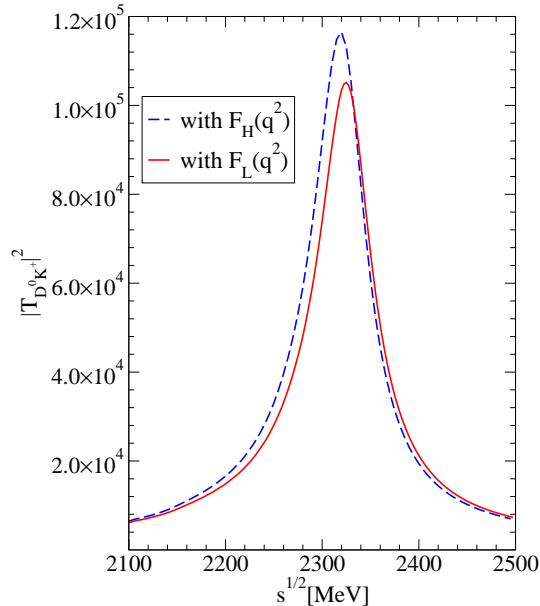
Figure 12: (Color online) Comparison of $|T|^2$ in the case of the $D_{s0}(2317)$ resonance for the two different form factors at $\rho = \rho_0$: type 1 (dashed line) with $F_H(q^2)$ of Eq. (26) and type 2 (solid line) with $F_L(q^2)$ of Eq. (19).

widths: the transparency ratio. The direct measurement in experiments of the in medium increased width is not easy because the observed decay channels usually come from the resonances that have escaped from the nucleus, so the density at the decay place is zero or very small [2]. Therefore, one should look at the production rate as a function of the mass number normalized to a particular nucleus (transparency ratio). This magnitude, which measures the survival probability, is very sensitive to the absorption rate of the resonance inside the nucleus, i.e., the in medium resonance width. This procedure has been succesfully used for the ϕ and ω production in nuclei in [4, 94] with the help of relatively easy tools of analysis [3, 95].

6   Conclusions

We have evaluated the selfenergy of low lying scalar mesons with open and hidden charm in a nuclear medium, concretely of the $D_{s0}(2317)$ and the theoretical hidden charm state $X(3700)$. The many body calculation has been done following the lines of previous studies in the renormalization of the light scalar mesons in the nuclear medium. The medium effects for the $D_{s0}(2317)$ and $X(3700)$ resonances are spectacular. Those resonances, which have zero and small width in free space, respectively, develop widths of the order of 100 and 200 MeV at normal nuclear matter density, respectively. The study also allowed us to trace back the reactions in the medium, which are responsible for the decay width of these mesons and which could be investigated in future reactions at hadron facilities. The option of looking at transparency

ratios was also suggested as a mean to investigate the widths of these mesons in nuclei. It was also discussed that the experimental study of this width and the medium reactions contributing to it provide information on the basic features of the resonance and the selfenergy of the D meson in a nuclear medium. In other words, the experimental analysis of those properties is a valuable test of the dynamics of the D meson interaction with nucleons and nuclei, and the nature of the charm and hidden charm scalar resonances, all of them topics which are subject of much debate at present. The results obtained here should stimulate experimental work in hadron facilities, in particular at FAIR [96], where the investigation of charm physics is one of the priorities.

Acknowledgments

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